

Supersymmetry Breaking and Gauge Mediation

RYUICHIRO KITANO¹, HIROSI OOGURI^{2,3} AND YUTAKA OOKOUCHI⁴

¹*Department of Physics, Tohoku University, Sendai 980-8578, Japan*

²*California Institute of Technology, Pasadena, CA 91125, USA*

³*Institute for the Physics and Mathematics of the Universe,
University of Tokyo, Kashiwa, Chiba 277-8582, Japan*

⁴*Perimeter Institute for Theoretical Physics, ON N2L2Y5, Canada*

Abstract We review recent works on supersymmetry breaking and gauge mediation. We survey our current understanding of dynamical supersymmetry breaking mechanisms and describe new model building tools using duality, meta-stability, and stringy construction. We discuss phenomenological constraints and their solutions, paying attentions to issues with gaugino masses and electroweak symmetry breaking.

CONTENTS

Introduction	2
Prototype scenario of low energy supersymmetry	5
Supersymmetry breaking	6
<i>Polonyi and generalized O’Raifeartaigh models</i>	<i>7</i>

<i>ISS model</i>	10
Gauge Mediation	11
<i>Analytic Continuation into Superspace</i>	12
<i>General Gauge Mediation</i>	13
Direct Mediation	16
<i>Breaking of R-symmetry</i>	17
<i>Anomalously small gaugino mass</i>	18
<i>Example of direct mediation model</i>	20
Stringy Realization	22
<i>Brane configuration</i>	22
<i>Holographic Gauge Mediation</i>	24
<i>Geometric Transition</i>	26
Electroweak symmetry breaking	27
NUMBERED LITERATURE CITED	31

1 Introduction

Supersymmetry has been playing important roles in modern particle physics even though there has been no direct experimental evidence for its existence in Nature. It is a hypothetical space-time symmetry which transforms bosonic states into fermionic ones and vice-versa, but this funny symmetry has nice features, such as vacuum stability ($E \geq 0$) and mild ultraviolet behavior of theories, *i.e.*, a restricted form of divergences. Many theorists expect that supersymmetry is an essential ingredient for the ultimate unified theory of elementary particles including gravity, perhaps the string theory.

Not only as a possible symmetry of Nature, supersymmetry have provided examples of (partly) calculable strongly coupled theories. Of the most amazing is the discoveries of strong/weak dualities among supersymmetric theories, such as dualities in the $\mathcal{N} = 4$ supersymmetric Yang-Mills theory (1,2,3,4,5), the Seiberg-Witten theory (6), the Seiberg dualities (7,8) in QCD-like supersymmetric gauge theories, and the gauge/gravity correspondence (9). These dualities boosted our understanding of non-perturbative physics.

With or without the connection to the theory of quantum gravity, particle theorists have applied supersymmetry to the Standard Model of Particle Physics. Their main motivation is to protect the Higgs potential against quantum corrections. Because of very weak ultraviolet divergences, one can naturally push the cut-off scale to an arbitrary high energy scale, such as the Planck scale ($\sim 10^{18}$ GeV).

Once one postpones the cut-off scale until, say, the Planck scale, it offers a wonderful arena for model-builders. One can build calculable field theoretical models of high-scale or high-temperature phenomena such as the inflation, baryogenesis, neutrino masses, fermion mass hierarchies, grand unification etc., without worrying about naturalness problems caused by large quantum corrections.

The successful unification of the coupling constants in the minimal supersymmetric extension of the Standard Model (10,11,12) has attracted many theorists to the supersymmetric world. Moreover, the lightest supersymmetric particle is a strong candidate for dark matter of the universe. See (13) for a review. Because of this beautiful and successful framework, there have been many studies on supersymmetry searches at high-energy colliders such as LEP, Tevatron, LHC, and ILC, and also through dark matter detections (14).

However, there is one missing piece, that is supersymmetry breaking mechanism. We need mass splittings between bosons and fermions because of experimental constraints. Satisfying the experimental bounds themselves is not hard. One can assume a supersymmetry breaking sector which couples to the minimal supersymmetric standard model (MSSM) in such a way that large enough mass splittings are obtained. A non-trivial constraint on model building comes from consideration of the Higgs sector. Naturalness of electroweak symmetry breaking together with the experimental bounds require that all the superpartners have masses of $\mathcal{O}(100 \text{ GeV})$.

In this article, we review recent developments towards a viable supersymmetric model of particle physics. We start with discussion of an early attempt at supersymmetric model building by Dine, Fischler, and Srednicki (15), in which the concepts of dynamical supersymmetry breaking, (direct) gauge mediation and their connection to the electroweak symmetry breaking have been discussed. Even though the specific model presented in the paper has been found to be incomplete due to progress in supersymmetric gauge theories, their scenario remains one of the most attractive and elegant. We then jump to recent progress and current understanding of those topics, including the discovery of meta-stable vacua in the supersymmetric QCD (SQCD) by Intriligator, Seiberg and Shih (ISS) (16), new formulations of gauge mediation and related topics, direct gauge mediation models and discussion on the gaugino masses, and connections to string theory. We close the discussion with open problems associated with electroweak symmetry breaking¹.

¹There are numbers of topics which we do not cover in this review article. Especially, we do not discuss cosmological or astrophysical constraints on supersymmetric models. We refer to, *e.g.*, (17, 18, 19) for constraints relevant for dynamical supersymmetry breaking and gauge

2 Prototype scenario of low energy supersymmetry

Witten in 1981 started asking a question of whether supersymmetry breaking happens dynamically. In (20), he proposed a natural framework for electroweak physics, “supersymmetric technicolor,” where a technicolor force dynamically breaks both supersymmetry and the electroweak $SU(2) \times U(1)$. Independently, in the same year, Dine, Fischler, and Srednicki (15), and Dimopoulos and Raby (21) proposed concrete models along the similar line. Especially in (15), a mechanism to generate gaugino and sfermion masses has been discussed. This mechanism is now called direct gauge mediation. We first recall this early attempt in this section.

Their picture is quite simple. There is a QCD-like $SU(M)$ gauge theory which becomes strong at a scale Λ_{SC} (SC stands for supercolor), and it is assumed that supersymmetry is dynamically broken by a condensation of a pair of the “superquarks,” ψ_S and $\bar{\psi}_S$, which are fermion components of chiral superfields, S and \bar{S} (Figure 1). Since a condensation of the fermion pair, $\langle \bar{\psi}_S \psi_S \rangle \simeq \Lambda_{\text{SC}}^3$, is a F -component of the meson superfield ($M \sim S\bar{S}$), supersymmetry will be spontaneously broken if such a condensation forms.

A part of the global symmetry in the sector is identified with the standard model gauge group by assigning quantum numbers to the superquarks. (In the original model only the $SU(2) \times U(1)$ part was embedded.) The gauginos and sfermions can obtain masses through loop diagrams involving the standard model gauge interactions (Figures 2 and 3). To generate non-zero gaugino masses, it is assumed that there is a boson-pair condensation

$$\langle \bar{s}s \rangle \simeq \Lambda_{\text{SC}}^2, \quad (1)$$

mediation.

in addition to the fermion-pair condensation. Here \bar{s} and s are the lowest components of \bar{S} and S , respectively. The presence of both condensations ensures breakdown of an R-symmetry necessary for non-vanishing gaugino masses. In order to obtain $\mathcal{O}(100)$ GeV gaugino and sfermion masses, the dynamical scale Λ_{SC} was assumed to be $\mathcal{O}(10)$ TeV due to the loop factors ($\sim \alpha/4\pi$).

Finally, the electroweak symmetry breaking can happen either through some other dynamics (technicolor) or the standard Higgs mechanism. In both cases, one can introduce elementary Higgs fields to write down the Yukawa interactions while avoiding the naturalness issue thanks to supersymmetry. In (15), a technicolor model is proposed where the electroweak VEV and the higgsino mass are generated through a strong dynamics at the $\mathcal{O}(300 \text{ GeV})$ energy scale as illustrated in Figures 4 and 5. (The possibility of using the usual Higgs mechanism was also discussed in the concluding section.)

We can see three essential ingredients in this model: dynamical supersymmetry breaking with R-symmetry violation, gauge mediation, and electroweak symmetry breaking through supersymmetry breaking. In the following, we will review recent progress and current understanding of those three components.

3 Supersymmetry breaking

In our current understanding, the assumption made in (15) for supersymmetry breaking,

$$\langle \bar{\psi}_S \psi_S \rangle \simeq \Lambda_{\text{SC}}^3, \quad (2)$$

is not valid in supersymmetric QCD theories. It has been shown that these theories have stable supersymmetric vacua (22, 23). Alternative possibilities have been considered and successful models for dynamical supersymmetry breaking

have been found, for example chiral gauge theories in (24, 25, 26) and theories with gauge singlet fields in (27, 28). See (29, 30) for more recent proposals.

Recently, dynamical supersymmetry breaking in supersymmetric QCD theories has revived by the work of ISS. They found a meta-stable supersymmetry breaking vacuum in $SU(N_c)$ supersymmetric gauge theories with massive N_f flavors for $N_c < N_f < 3N_c/2$. The presence of the vacuum is established when $m \ll \Lambda$ with m and Λ the quark mass and the dynamical scale, respectively. A non-vanishing fermion-pair condensation is obtained to be

$$\langle \bar{\psi}_Q \psi_Q \rangle \sim m \Lambda^2, \quad (3)$$

instead of Equation (2) as illustrated in Figure 6.

Let us first discuss the linear sigma models for supersymmetry breaking which are widely used as a tool to establish existence of supersymmetry breaking vacuum in strongly coupled theories (24, 25, 26, 31, 32) and also as effective theories to describe low energy physics of a variety of supersymmetry breaking models. We then introduce the ISS model and its connection to the linear sigma model. For a recent pedagogical review on supersymmetry breaking, see (33).

3.1 Polonyi and generalized O’Raifeartaigh models

Supersymmetry is spontaneously broken when the F -component of a chiral superfield X acquires a VEV:

$$\langle F_X \rangle \neq 0, \quad (4)$$

since $2\langle F_X \rangle = \langle \{iQ, X\} |_{\theta} \rangle$ which should vanish if $Q|0\rangle = 0$. It is easy to construct a linear sigma model to describe this phenomenon using a single chiral superfield

X . The superpotential is simply,

$$W = \mu^2 X. \quad (5)$$

The F -component of X acquires a VEV unless the Kähler potential is singular. This is the unique choice of the superpotential for a model with a single chiral superfield X . (There can be a small perturbation to it if one allows the vacuum to be meta-stable.) By an appropriate shift of X we can choose the stable point at $X = 0$. In order for this point to be stable, the Kähler potential expanded around it is assumed to be of the form:

$$K = X^\dagger X - \frac{(X^\dagger X)^2}{\Lambda^2} + \dots \quad (6)$$

Equations (5) and (6) define the linear sigma model for supersymmetry breaking, the Polonyi model (34).

The equation of motion for F_X and the potential minimization of X leads

$$\langle F_X \rangle = \mu^2, \quad \langle X \rangle = 0. \quad (7)$$

Therefore, supersymmetry is spontaneously broken. The fermion component ψ_X remains massless. This is the Goldstino fermion associated with the spontaneous supersymmetry breaking. The complex scalar field X obtains a mass:

$$m_X^2 = \frac{4\mu^4}{\Lambda^2}. \quad (8)$$

Up to $\mathcal{O}(1/\Lambda^2)$ there is no mass splitting between the scalar and the pseudo-scalar parts due to an unbroken approximate R-symmetry with $R(X) = 2$.

Although it is a non-renormalizable model, the model serves as the effective theory (if $\mu^2 \ll \Lambda^2$) for a wide class of supersymmetry breaking models. Conversely, one can establish the presence of a supersymmetry breaking vacuum when a model reduces to the Polonyi model at a point of the field space

by integrating out massive degrees of freedom. This technique has been used in (24, 25, 26, 31, 32, 16, 35, 36, 37).

One can also use the Polonyi model as a hidden sector whose supersymmetry breaking effects are communicated to the MSSM sector through some interactions such as gravity or gauge interactions. Especially in (38, 39, 40, 41), simple models of gauge mediation have been constructed by attaching the messenger fields to the above linear sigma model. They belong to the indirect type of gauge mediation in contrast to the direct one depicted in Figure 2 and discussed later.

The O’Raifeartaigh model (42) and its generalization are alternative simple theories for supersymmetry breaking. (They reduce to the Polonyi model when massive fields are much heavier than the size of supersymmetry breaking.) The generalized O’Raifeartaigh model is defined by a set of chiral superfields (Φ_i) and a superpotential up to cubic terms (the Wess-Zumino model (43)) where

$$\frac{\partial W}{\partial \Phi_i} = 0 \tag{9}$$

cannot be solved simultaneously. The chiral superfields Φ_i have the canonical Kähler potential, $K = \Phi_i^\dagger \Phi_i$. These models have received renewed attention recently because many dynamical models (including the ISS model) reduce to it at low energy. One of the basic property common to all the O’Raifeartaigh models is existence of tree-level flat directions, known as pseudo-moduli, emanating from any local supersymmetry breaking vacuum (44). Such pseudo-moduli receive quantum corrections and a potential along the flat direction is generated.

3.2 ISS model

It has been known from the studies by Seiberg (8) that supersymmetric QCD theories with the quark mass m has stable supersymmetric vacua at

$$\langle M \rangle = \langle Q\bar{Q} \rangle \sim \Lambda^{(3N_c - N_f)/N_c} m^{(N_f - N_c)/N_c}, \quad (10)$$

where N_c and N_f are numbers of color and flavor, respectively. This fact appeared to be a no-go theorem for dynamical supersymmetry breaking in these theories (22).

Recently, it was shown by ISS that there can be a meta-stable supersymmetry breaking vacuum far away from the supersymmetric vacua. The model they considered is a supersymmetric $SU(N_c)$ theory with N_f fundamental and anti-fundamental chiral superfields Q_i and \bar{Q}_i ($i = 1, \dots, N_f$). The superpotential is

$$W = m Q_i \bar{Q}_i, \quad (11)$$

where the color and flavor indices are contracted. This is an asymptotically free theory and becomes strongly coupled in the IR. Therefore, the analysis in the electric picture (description in terms of quarks and gluons) is difficult near the origin of the meson space $M \sim Q\bar{Q}$. By the power of duality, the IR physics of the model is described by the dual magnetic theory, which is weakly coupled for $N_c < N_f < 3N_c/2$ (7). Perturbative calculations in the magnetic theory are reliable near the origin of the meson field space.

The dual gauge theory is $SU(N_f - N_c)$ with N_f dual quarks, q and \bar{q} , and the gauge singlet meson fields M . The superpotential is

$$W_{\text{mag.}} = h\mu^2 M_{ii} - h q_i M_{ij} \bar{q}_j, \quad (12)$$

where h is a coupling constant of $\mathcal{O}(1)$ and the term $h\mu^2 M$ corresponds to the quark mass terms in Equation (11). The coefficient $h\mu^2$ is naturally of $\mathcal{O}(m\Lambda)$. Again, i and j are the $SU(N_f)$ flavor indices. At the point where $\langle M_{ij} \rangle = 0$ and $\langle q_i \rangle = \langle \bar{q}_i \rangle \neq 0$ and the dual gauge group is completely broken, the flavor $SU(N_f)$ symmetry is broken down to $SU(N_f - N_c) \times SU(N_c)$. The effective superpotential at that point has the form of the O’Raifeartaigh model²:

$$W = h\text{Tr} \left(\mu^2 X - X\rho\tilde{\rho} - \mu\rho\tilde{Z} - \mu Z\tilde{\rho} \right),$$

where X is the pseudo-moduli field, which is a $N_c \times N_c$ part of the meson M , and ρ , $\tilde{\rho}$, Z and \tilde{Z} are massive fields. Once we integrate out these massive fields, the Coleman-Weinberg potential is generated, which stabilize the pseudo-moduli at the origin (16). The F -component of the X field ($\sim \bar{\psi}_Q \psi_Q$) gets a VEV triggered by an explicit breaking of chiral symmetry as in Figure 6. Although the detailed structure is different, the dynamical supersymmetry breaking through a fermion-pair condensation assumed in (20, 15, 21) is revived in the meta-stable vacua of supersymmetric QCD models.

The tunneling rate into the true supersymmetric vacuum in Equation (10) is exponentially suppressed if $m \ll \Lambda$ since the supersymmetric vacuum is located far from the meta-stable vacuum, in comparison with the height of the potential along the meson direction $V^{1/4} \sim (m^2 M^2)^{1/4}$.

4 Gauge Mediation

Inspired by the simplicity and genericity of meta-stable supersymmetry breaking, model building of gauge mediation have been invigorated lately³. The diagram in

²For simplicity we omitted other pseudo-moduli in the model. For detail, see (16).

³See (45) for a review on earlier works on gauge mediation.

Figure 2 has been reconsidered with new knowledge of supersymmetry breaking, and also been reformulated by using current correlators. In this section, we review techniques to calculate the gaugino and sfermion masses.

4.1 Analytic Continuation into Superspace

One can perform explicit calculations of gaugino and sfermion masses when the supercolor boxes in Figures 2 and 3 are loops of weakly coupled fields, called the messenger fields. Such a class of models are first considered in (46, 47, 48) and revived in mid '90s by (49, 50, 51) where explicit computations of gaugino and sfermion masses are performed. The highly predictive feature of gauge mediation is demonstrated; the gaugino and sfermion masses are functions only of their quantum numbers to a good approximation.

In (52, 53), a powerful method was developed to compute multi-loop quantities from one-loop running data. The results follow from imposing constraints due to holomorphy in a spurion $X = X_0 + \theta^2 F$ on the effective action. One can use the method as the leading order calculation in the F/X_0^2 expansion; higher order terms in F arise from terms involving super-derivatives, which are not considered. The gaugino masses are identified with the F -component of the holomorphic gauge coupling, $\tau(X)$. The X dependence is originated from the change of the beta functions due to the decoupling of a messenger field whose mass is X_0 in the supersymmetric limit. The sfermion masses can be extracted from the wave-function renormalizations $Z_Q(X, \bar{X})$ of the matter superfields Q . At the leading order, Z_Q depends on X through the gauge couplings $\tau(X)$. By expanding the effective action in F , one can describe the gaugino and sfermion

masses by one-loop quantities,

$$m_\lambda = -F \left. \frac{\partial \ln \tau(X)}{\partial X} \right|_{X=X_0} \simeq \Delta\beta_{g_a}^{(1)} \frac{F}{X_0},$$

$$m_s^2 = -|F|^2 \left. \frac{\partial^2 \ln Z_Q(X, X^\dagger)}{\partial X \partial X^\dagger} \right|_{X=X_0} \simeq \gamma_s^{(1)} \Delta\beta_{g_a}^{(1)} \left| \frac{F}{X_0} \right|^2, \quad (13)$$

where $\Delta\beta_{g_a}^{(1)}$ is a discontinuity of the coefficient of the one-loop beta function at the threshold X_0 , and $\gamma_s^{(1)}$ is one-loop anomalous dimension of the matter fields.

This method simplifies calculations of soft masses.

In addition to applications for gauge mediation, the technique of the analytic continuation into superspace can be used in computing the effective potential of the pseudo-moduli in the supersymmetry breaking sector. In (54,55) a leading-log effective potential for a pseudo-moduli is computed. It turned out that, even if the potential is generated only at higher loops, there is a regime where the potential can be simply determined from a combination of one-loop running data. The results were applied to survey pseudo-moduli spaces for large classes of models (55).

4.2 General Gauge Mediation

The method in subsection 4.1 only works in cases where the messenger fields are weakly coupled. In general, in order to evaluate diagrams such as in Figures 2 and 3, we need information on two-point functions of currents in the supersymmetry breaking dynamics.

Recently the authors of (56) provided the most general parametrization of the two-point functions by giving a model-independent definition of gauge mediation: In the limit that the MSSM gauge coupling $\alpha_i \rightarrow 0$, the MSSM sector decouples from the hidden sector that breaks supersymmetry. In particular, the MSSM

gauge group becomes a global symmetry G of the hidden sector in this decoupling limit. All the information we need for gauge mediation is encoded in the currents and their correlation functions.

The conserved currents are real linear supermultiplets satisfying $D^2 \mathcal{J} = 0$. For a $U(1)$ symmetry, the current superfield can be written in components as

$$\mathcal{J} = J + i\theta j - i\bar{\theta}\bar{j} - \theta\sigma^\mu\bar{\theta}j_\mu + \dots$$

Two-point functions are constrained by the Lorentz invariance and the current conservation, as follows

$$\langle J(p)J(-p) \rangle = \tilde{C}_0(p^2/M^2),$$

$$\langle j_\alpha(p)\bar{j}_{\dot{\alpha}}(-p) \rangle = -\sigma_{\alpha\dot{\alpha}}^\mu p_\mu \tilde{C}_{1/2}(p^2/M^2),$$

$$\langle j_\mu(p)j_\nu(-p) \rangle = -(p^2\eta_{\mu\nu} - p_\mu p_\nu)\tilde{C}_1(p^2/M^2),$$

$$\langle j_\alpha(p)j_\beta(-p) \rangle = \epsilon_{\alpha\beta}M\tilde{B}(p^2/M^2).$$

Gaugino and scalar masses at leading order in the gauge coupling are governed by the two-point functions⁴,

$$m_\lambda = g^2 M \tilde{B}(0), \quad m_s^2 = g^4 Y^2 A, \quad (14)$$

where Y is the $U(1)$ charge of the sfermion and A is defined by

$$A = - \int \frac{d^4 p}{(2\pi)^4} \frac{1}{p^2} \left(3\tilde{C}_1(p^2/M^2) - 4\tilde{C}_{1/2}(p^2/M^2) + \tilde{C}_0(p^2/M^2) \right).$$

Note that A is a linear combination of terms with different signs, thus sfermion masses squared are not necessarily positive. Even under this generic situation it

⁴In general one-point function $\langle J \rangle$ is non-zero, which could lead to tachyonic sfermion. To forbid a contribution from the one-point function, the messenger \mathbf{Z}_2 parity was imposed $\mathcal{J} \rightarrow -\mathcal{J}$ in (56).

was derived that the scalar masses obey two sum rules,

$$\text{Tr}(Y m_s^2) = 0, \quad \text{Tr}((B - L) m_s^2) = 0.$$

In generalizing the $U(1)$ theory to full Standard Model gauge group $SU(3) \times SU(2) \times U(1)$, A and B are replaced by A_k and B_k ($k = 1, 2, 3$) corresponding to the three gauge group factors. These are three real and complex numbers, providing nine independent parameters in general gauge mediation. Arbitrary phases of B_k would typically lead to an unacceptable level of the electric dipole moment of an electron and a neutron. So it is plausible to assume that the hidden sector is CP invariant and to impose phenomenological constraints on the phases. Thus we are left with six real parameters that span the parameter space of the general gauge mediation.

It is interesting to ask if there are simple models of weakly coupled messengers that cover the entire parameter space. This question was first raised in (57) and studied in F -term supersymmetry breaking models. The authors found models with the right number of parameters, but these models have turned out not to cover the entire parameter space. This question was examined further in (58), using D -term supersymmetry as suggested in (59). The authors found models with weakly coupled messengers that span the whole parameter space. Though direct mediations include cases with strongly coupled hidden sectors, the answer posed at the beginning of this paragraph has turned out to be positive.

The idea of analytic continuation into superspace reviewed in subsection 4.1 was extended to general gauge mediation for small supersymmetry breaking scale, $F \ll M^2$ (60,61). The authors of (60,61) showed identities for the two-point functions and reproduced the result shown in Equation (13) by exploiting the general mass formulae, Equation (14). For more works on general gauge mediation, see

(62, 63, 64, 65, 66).

5 Direct Mediation

With the technology of calculating or parametrizing the gaugino/sfermion masses, we can proceed to model building with direct gauge mediation. Direct mediation is a class of gauge mediation where messenger sectors are responsible for supersymmetry breaking and (meta) stability of vacua⁵, *e.g.*, Figures 2 and 3 as the ultimate version of it. In some cases, supersymmetry is restored if couplings to messengers are turned off. For a more precise definition of direct mediation, see (71, 57). Following the seminal work (15, 21), this possibility of model building has received much attention. See (72, 73, 74, 75, 76) for some of the earlier works.

In this section, we will discuss two major challenges in constructing phenomenologically viable models and their possible solutions: the Landau pole and the light gaugino mass problems.

In a direct-type model, the messenger sector cannot be adjusted arbitrarily since it is closely tied to supersymmetry breaking effects. Especially, if there are a large number of the messenger fields, the coupling constants of the standard model gauge interaction hit the Landau pole below the unification scale. One can see difficulties for example in the models of (77, 78, 79). In (80), a model was proposed to alleviate this issue by using a non-trivial conformal fixed point. For a more recent progress in this approach, see (81).

Another issue is how to generate large enough gaugino masses. Non-zero Ma-

⁵There is also semi-direct mediation, where messengers couple to the supersymmetry breaking sector but are not relevant for the stability of a vacuum (67). It includes mediator models considered in (53). Such models have been studied further in (68, 69, 70).

jorana gaugino masses require an R-symmetry to be broken⁶. However, simple O’Raifeartaigh type models including the ISS model preserve an (approximate) R-symmetry at the supersymmetry breaking vacuum. One needs a careful arrangement to achieve supersymmetry breaking without R-symmetry. In addition, R-symmetry breaking is often not sufficient to guarantee large enough gaugino masses. Below we will review these issues and discuss how to construct successful models.

5.1 Breaking of R-symmetry

An R-symmetry can be broken either spontaneously or explicitly. In (88), it was shown that, for a generalized O’Raifeartaigh model⁷, spontaneous R-symmetry breaking requires a field with an R-charge different from 0 or 2. Explicit examples of direct mediation with exotic R-charges were constructed in (90). Since such *extra ordinary* gauge mediation models do not necessarily preserve the approximate messenger parity (91), one-loop corrections generate tachyonic scalar masses in general. A way to suppress such one-loop contributions was discussed in (57).

Another possibility is to break an R-symmetry explicitly at the Lagrangian level, which has an added advantage of avoiding the unwanted R-axion. According to (23), for a generic superpotential, stable supersymmetry breaking vacua

⁶An alternative approach is to use Dirac gaugino mass (82, 83, 84, 85). This requires extra fields in adjoint representation in the MSSM, and such models typically suffer from the Landau pole problem. It is interesting to note that the R-symmetry does not have to be broken in such models. See (86, 87) for recent proposals.

⁷This statement does not hold for a model with gauge interactions (71) or a model in which two-loop Coleman-Weinberg potential is dominant (89).

require an unbroken R-symmetry. However, explicit R-symmetry breaking is allowed if we can live with meta-stable vacua. In particular, if R-symmetry breaking is small, supersymmetric vacua are generated near the infinity in the field space, which guarantees the meta-stability of the supersymmetry breaking vacuum. Since the R-symmetry is broken, there is no symmetry reason to prohibit the gaugino masses. However, it turns out that generating large enough gaugino masses is still a challenge, as we will explain now.

5.2 Anomalously small gaugino mass

Let us consider the case where the supersymmetry breaking scale \sqrt{F} is much smaller than the messenger scale M . The leading contribution to gaugino masses in the F/M^2 expansion is given by

$$m_\lambda = \frac{g^2}{(4\pi)^2} \cdot F_X \frac{\partial}{\partial X} \log[\det \mathcal{M}(X, M)], \quad (15)$$

where $\mathcal{M}(M)$ is fermion mass matrix of messengers and X is the pseudo-moduli field with non-zero F -component responsible for supersymmetry breaking (52). The right-hand side of this equation vanishes in a simple O’Raifeartaigh model even if the R-symmetry breaking is large (92). This problem has been observed quite often in models of direct gauge mediation (71, 77, 93, 78, 94). Therefore, it was forced to consider models with $F/M^2 = \mathcal{O}(1)$, *i.e.*, $\sqrt{F} \sim M \sim 10 - 100$ TeV, with which it is difficult to avoid the Landau pole problem. Moreover, it is interesting that the models with vanishing leading order gaugino masses are severely constrained by a recent Tevatron bound on the sparticle masses and a mass bound on a light gravitino (95).

Recently, the reason for the small gaugino masses was explained in (96) on a general ground. Suppose that the low energy effective theory near a given vacuum

is described by the O’Raifeartaigh model with the superpotential,

$$W = fX + \frac{1}{2}(\lambda_{ab}X + m_{ab})\phi_a\phi_b + \frac{1}{6}g_{abc}\phi_a\phi_b\phi_c. \quad (16)$$

By performing a unitary transformation of fields, one can rewrite the general O’Raifeartaigh model in this form (44). The fields ϕ_a are identified with messengers. The Kähler potential is canonical for all the chiral superfields. In this model, it is proven that $\det(\lambda X + m)$ is a constant (independent of X) in a *stable* supersymmetry breaking vacuum, *i.e.*, in the lowest energy state (96). In the proof, it is assumed that there is no unstable point anywhere in the pseudo-moduli space since otherwise there should be a lower energy state.

This theorem immediately means that the leading-order gaugino masses vanish at the leading order in the F/M^2 expansion in this vacuum (from Equation (15) with $\mathcal{M} = \lambda X + m$), regardless of whether or how an R-symmetry is broken. Since many dynamical supersymmetry breaking models are effectively described by Equation (16) near the vacuum, this puts constraints on model building with direct gauge mediation. The meta-stability does not help qualitatively since the presence of a supersymmetric vacuum should be a small effect. It is interesting that anomalously small gaugino masses is related to global structure of the vacua of the theory and does not depend on details on how the R-symmetry is broken.

One way to solve this gaugino mass problem is to choose our vacuum to be even more meta-stable, *i.e.*, to use an even higher energy state compared to the lowest supersymmetry breaking vacuum so that the presence of an unstable point in the pseudo-moduli space is allowed. We will now describe such a model, which antedated the theorem in (96).

5.3 Example of direct mediation model

As we have discussed already, a successful model for large enough gaugino masses requires R-symmetry breaking as well as a certain global structure of the pseudo-moduli space. We review in this subsection the model of (80) as an example of models with non-vanishing gaugino masses at the leading order in F . It is a deformation of the ISS model in section 3: an $SU(N_c)$ gauge theory with N_f flavors with mass terms.

The idea is to embed the standard model gauge group into the global symmetry of the ISS model, $SU(N_f) \times U(1)_B$. In order to guarantee (meta)stability, to avoid the Landau pole below the unification scale, and to generate non-vanishing gaugino masses, it is assumed that,

- (1) the quark masses are split into two groups; m_{light} for N_c flavors (Q_a , $a = 1, \dots, N_c$) and m_{heavy} for $N_f - N_c$ flavors (Q_I , $I = 1, \dots, N_f - N_c$), and
- (2) the presence of a quartic term, $(Q_I \bar{Q}_a)(Q_a \bar{Q}_I)$, in the superpotential. The quartic term breaks the R-symmetry explicitly. It makes the structure of the pseudo-moduli space richer. As we will see, this helps to generate sizable gaugino masses.

The global symmetry is now $SU(N_f - N_c) \times SU(N_c) \times U(1)_B$. It is shown that the Landau pole can be avoided when the standard model gauge group is embedded to the $SU(N_f - N_c)$ part and $m_{\text{heavy}} \gg m_{\text{light}}$.

In the magnetic description, the effective superpotential is given by

$$W = h \mu^2 \sum_{a=1}^{N_c} M_{aa} + h m^2 \sum_{I=1}^{N_f - N_c} M_{II} - h \text{Tr } q M \bar{q} - h m_z \sum_{a,I} M_{Ia} M_{aI}, \quad (17)$$

where h is a parameter of $\mathcal{O}(1)$ and $\mu \ll m$ for $m_{\text{light}} \ll m_{\text{heavy}}$. The last term (the term with a coefficient $h m_z$) corresponds to the quartic term in the

electric description. A part of the meson fields, M_{II} , M_{aI} , and M_{Ia} , and of the dual quarks, q_I and \bar{q}_I , are charged under the standard model gauge group and thus are messenger fields. In addition to the perturbed ISS vacuum where the meson fields M_{aa} acquire small VEVs which vanish as $m_z \rightarrow 0$, there are new supersymmetry breaking vacua where the meson fields acquire large VEVs which go to infinity as $m_z \rightarrow 0$. The new vacua has lower energies than the perturbed ISS vacuum.

In the new vacua, gaugino masses were found to vanish at the leading order of F . We now understand the reason for it, thanks to the theorem of (96).

On the other hand, the original ISS vacuum has higher energy and thus evades the theorem. Near the origin of the meson space, the low energy effective theory is the O’Raifeartaigh model with the superpotential,

$$W = h\text{Tr} \left(\mu^2 X - X\rho\tilde{\rho} - me^{-\theta}\rho\tilde{Z} - me^{\theta}Z\tilde{\rho} - m_z Z\tilde{Z} \right).$$

One-loop effects stabilize the potential, giving rise to the ISS vacuum. However, the tree level scalar potential has an unstable region near the new vacua, $X \sim m^2/m_z$. It is because of this that the ISS vacuum evades the theorem of (96) and generates the gaugino masses,

$$m_\lambda = \frac{g^2 N_c}{(4\pi)^2} \frac{h\mu^2}{m} \frac{m_z}{m} + O\left(\frac{m_z^2}{m^2}\right), \quad (18)$$

where g is the coupling constant of the standard model gauge interaction. Scalar masses are also obtained by two-loop diagrams and can be adjusted to be the same order as the gaugino mass $\mathcal{O}(m_i) \simeq \mathcal{O}(m_\lambda)$ by setting $m_z \sim m/\sqrt{N_c}$.

6 Stringy Realization

We are going to have an interlude to describe realization of supersymmetry breaking mechanisms and their mediations in string theory. There are two motivations for stringy constructions. One is to provide a unified framework for field theory models to understand the nature of the parameters of these models. Moreover, string theory has been used to develop powerful computational tools for field theory effects and to gain geometric insights into supersymmetry breaking phenomena. Another motivation is to understand supersymmetry breaking in string theory better. In the past quarter-century, much progress has been made in understanding of string theory in supersymmetric backgrounds and many exact results on non-perturbative effects have been derived. It is desirable to extend these results to non-supersymmetric situations.

Once we accept meta-stable vacua, a wide variety of stringy realizations becomes possible. Since long-lived meta-stable vacua are ubiquitous in vector-like models such as SQCD, it is relatively easy to embed supersymmetry breaking models in string theories. In this section, we will review model building by intersecting branes, gauge/gravity dualities and geometric transitions. For an interesting recent development of string phenomenology with F -theory, see (97) for a review.

6.1 Brane configuration

Open strings are collective coordinates of D branes. On parallel Dp branes, if we take the limit of string length $\ell_s \rightarrow 0$ while scaling the string coupling constant as $g_s = g_{\text{YM}}^2 \ell_s^{3-p}$, the low energy effective theory becomes the maximally supersymmetric Yang-Mills theory in $(p+1)$ dimensions with the gauge coupling given by

g_{YM} . We can reduce the number of supercharges by considering configurations of intersecting branes. For example, if we have a pair of parallel NS 5-branes, and if we suspend parallel D4 branes between them in such a way that all the branes share 4 dimensions, then the low energy effective theory is the $\mathcal{N} = 2$ supersymmetric pure Yang-Mills theory in 4 dimensions. Preserving supersymmetry requires that branes be in specific relative angles. Any other angle would break supersymmetry completely. See (98) for a review.

Stringy description of the ISS model and its meta-stable vacua was initiated in (99,100) and studied further in (101,102). Figure (7) shows the brane configuration for $\mathcal{N} = 1$ $SU(N_c)$ gauge theory with N_f fundamental chiral multiplets with masses m_1, \dots, m_{N_f} , which is the ISS model. The meta-stable supersymmetry breaking vacua of the ISS model can be found by going to the magnetic description. To go from the electric description to the magnetic dual, we exchange the location of the two NS 5-branes. The number of branes changes during this process due to the Hanany-Witten mechanism. Figures (8.a) and (8.b) show the electric and magnetic branes configurations when masses are equal to zero. We see that the gauge group in the magnetic description is $SU(N_f - N_c)$ as expected by the Seiberg duality. In both description, D4 branes are parallel and supersymmetry is preserved.

Now, let us turn masses to non-zero. In the electric description, this can be done by moving the D6 branes up and down as in Figure (7). In the magnetic description, however, we cannot move the D6 branes in this way while keeping all the D4 branes parallel to each other. We can keep $(N_f - N_c)$ D4 branes in parallel by reconnecting them at the NS 5-branes, but the remaining N_c D4 branes have to be tilted as shown in Figure (9). The resulting configuration breaks all the

supersymmetry.

The brane configuration in this limit reproduces various features of vacuum structure of the ISS model such as global symmetries, pseudo-moduli and vacuum energy, as explained in (99,100). Thus, it is reasonable to identify this configuration with the supersymmetry breaking meta-stable vacuum in the field theory.

In (101), it was claimed that this brane configuration is not related to the original electric description and that it does not describe the ISS vacuum because of a certain asymptotic behavior of D branes caused by brane bending at finite g_s . In (102), an alternate interpretation of the brane bending involving string tachyons was proposed and it was argued that the brane configuration of (99,100) indeed describes the ISS vacuum. We refer to these papers for details on the issue.

6.2 Holographic Gauge Mediation

Soft-terms in gauge mediation models are related to correlators of global symmetry currents in the hidden sector (56), as we reviewed in section 4.2. Such relations are particularly useful for direct-type models where hidden sectors are strongly coupled but the current correlators may be calculable. In (63), this idea was applied to a class of strongly coupled hidden sectors to derive low energy parameters of the models. Gauge/gravity duality in string theory can be useful since the current correlators in field theories are computed by Green's functions of the corresponding gauge fields in the gravity duals.

One of the well-studied examples is the duality between the cascading $\mathcal{N} = 1$ $SU(N + M) \times SU(N)$ gauge theory and the warped deformed conifold geometry in type IIB string theory (103). In (104), a meta-stable supersymmetry breaking

vacuum was constructed in this setup by putting anti-D3 branes at the tip of this conifold geometry. This vacuum can decay into a supersymmetric vacuum by quantum tunneling and brane/flux annihilation. Since the process happens near the tip of conifold, the UV description is not modified and we expect that there is a field theory dual to describe the process. Both the supersymmetric vacuum and the meta-stable vacuum must have corresponding states in such a field theory. Recently, this construction was generalized in (105,106) to quiver gauge theories and string theory on quotients of the conifold, and the correspondence between meta-stable vacua in the field theories and in the gravity theories have been clarified. In particular, the model studied in (106) gives a stringy realization of the direct gauge mediation model constructed in (80) and discussed in the previous section. Moreover, it provides a natural mechanism to generate small parameters in the model.

In the holographic gauge mediation model⁸ of (112), strongly coupled messengers and hidden sectors are replaced by the supersymmetry breaking solutions of Type IIB supergravity constructed in (113), which takes into account the back reaction of the anti-branes. Although gaugino condensation breaks the R-symmetry to \mathbf{Z}_2 , the issue of anomalously small gaugino masses is not resolved in this model. This construction was generalized in (114), by using the supersymmetry breaking solution of (115).

⁸For earlier related works in five dimensional theory in a truncated AdS_5 background, see (107, 108, 109, 110, 111).

6.3 Geometric Transition

Another way to construct low energy gauge theories is to use D branes wrapping cycles of Calabi-Yau manifolds. For example, consider a resolved conifold and wrap its small 2-sphere with N D5 branes. The low energy effective theory is the $\mathcal{N} = 1$ pure Yang-Mills theory in 4 dimensions with gauge group $U(N)$. By a chain of dualities, one can relate this construction to an intersecting brane configuration of the type discussed earlier in this section (116).

The geometric transition relates the resolved conifold to a deformed conifold with a small 3-sphere. In particular, the Veneziano-Yankielowicz superpotential of the $\mathcal{N} = 1$ Yang-Mills theory realized on the N D5 branes on the resolved conifold is equal to the Gukov-Vafa-Witten superpotential for the deformed conifold with N units of Ramond-Ramond fluxes through the 3-sphere (117). This is an example of the large N duality and is closely related to the gauge/gravity duality of (103, 118). It has been generalized to a variety of gauge theories and corresponding geometries, and their superpotentials are computed exactly using topological string theory (119).

In (117), it was suggested that the geometric transition can be used even when both branes and anti-branes are present. This proposal was made more precise in (120, 121, 122). It was conjectured that the large N limit of brane/anti-brane systems are Calabi-Yau manifolds with fluxes and that the physical potential can be computed using the method of topological string theory as in supersymmetric cases, except that some of the fluxes are negative.

In (122, 123), field theory descriptions of such meta-stable vacua are given. Their M-theory realizations have been explored in (124, 125).

7 Electroweak symmetry breaking

The final topic is the most relevant for particle physics. In the MSSM, electroweak symmetry breaking does not happen in the supersymmetric limit since supersymmetry guarantees the absence of the negative mass squared for the Higgs boson. Thus, there must be a coupling between the supersymmetry breaking dynamics and the Higgs sector of the MSSM.

The use of a technicolor dynamics to break both supersymmetry and electroweak symmetry was proposed in (20,15,21). The Higgs fields obtain VEVs through a direct coupling to the dynamics, and give masses to fermions in the Standard Model. On the other hand, in calculable models such as in (51,50,49), one usually assumes that the Higgs sector is separated from the supersymmetry breaking dynamics and some communication through messenger fields generates the Higgs potential and the Higgsino mass.

There are, in fact, number of difficulties in this program and completely successful models have not been found. Here we explain the difficulties and current status. First, we list the phenomenological requirements:

1. The Z boson mass, $m_Z = 91$ GeV,
2. The Higgs boson mass constraint, $m_h > 114$ GeV (126),
3. The Higgsino mass constraint, $m_{\tilde{H}} > 94$ GeV (127),
4. The gaugino mass constraint, $m_\lambda > O(100)$ GeV (127)⁹,
5. The top-quark mass, $m_t = 173$ GeV (131),
6. The bound on the electric dipole moments of the neutron, $d_n < 3 \times 10^{-26} e$ cm (132).

⁹One of the neutralinos can be very light (128,129,130).

As we will explain below, it is not easy to naturally explain all of the above in supersymmetric models. This is generally called the μ -problem. Most of the discussion below are not specific to gauge mediation models, but the problems are particularly sharp in gauge mediation since it is designed to be calculable.

The Higgs potential in the MSSM is

$$V = (m_{H_u}^2 + \mu^2)|H_u|^2 + (m_{H_d}^2 + \mu^2)|H_d|^2 + (B\mu H_u H_d + \text{h.c.}) + \frac{1}{8}(g_Y^2 + g_2^2)(|H_u|^2 - |H_d|^2)^2, \quad (19)$$

where H_u and H_d are the neutral components of the Higgs fields in the MSSM. The part of the quadratic terms μ^2 is supersymmetric contribution to the Higgs potential. By definition, μ is the mass of the Higgsino. By supersymmetry, the quartic potential is related to the gauge couplings, g_Y and g_2 , which control the Higgs boson mass. Other quadratic terms are soft supersymmetry breaking terms which should arise from couplings to the supersymmetry breaking sector. The B parameter is in general complex valued.

If the Higgs fields are only weakly coupled to the supersymmetry breaking sector such as in (15), the typical diagram to generate the Higgsino mass (the μ -term) is given in Figure 5 where the technicolor box is replaced by a messenger loop or a supersymmetry breaking box. The problem is that if there is a one-loop diagram as in Figure 5, *one-loop* diagrams for generating $m_{H_u}^2$, $m_{H_d}^2$ and $B\mu$ should also be present. Since μ^2 is effectively *two-loop* valued, it suggests that the typical mass scale for the Higgs potential is larger than the Higgsino mass by a one-loop factor, *i.e.*, $m_Z \gg \mu$. Now we can see inconsistency with the items 1 and 3.

One could have evaded the above problem by assuming that the Higgs fields do not acquire VEVs from their potential and there are other sources for the Z

boson mass such as technicolor models in (15, 21). Small VEVs can be obtained through the diagram in the right-hand-side of Figure 4, and they are responsible for the fermion masses. However, since the top quark is rather heavy (item 5), such an assumption is not easily realized in a consistent way.

Even if we somehow avoid the hierarchies in the parameters in the Higgs potential at the messenger scale dynamics, there are also one-loop corrections to $m_{H_u}^2$ and $m_{H_d}^2$ which are proportional to the gaugino masses squared. Therefore, the gaugino masses cannot be very large compared to the Z -boson mass. Together with the item 4, the gaugino masses are $\mathcal{O}(100)$ GeV. In gauge mediation, the gaugino masses are obtained at one-loop level. Therefore, it suggests that the Higgs fields are also weakly coupled to the supersymmetry breaking sector. In this case, the hierarchy mentioned above among $m_{H_u}^2$, $m_{H_d}^2$, $B\mu$ and μ^2 is a quite generic consequence.

In the Higgs potential in Equation (19), there is a tree-level prediction for the Higgs boson mass: $m_h \leq m_Z$, which is clearly inconsistent with the items 1 and 2. This requires a rather large quantum correction to the quartic couplings arising from supersymmetry breaking (133, 134, 135). In the MSSM, the largest contribution is from the top-stop loop diagrams. However, if there is a large quantum correction to the quartic coupling constant, there is also a large contribution to the quadratic term, especially to $m_{H_u}^2$. Such a large contribution calls for fine-tuning of the parameters to have a correct size of the electroweak VEV, *i.e.*, the Z -boson mass (136) (See (137, 138) for recent discussions.).

The degree of fine-tuning gets relatively milder when we have a sizable stop-stop-Higgs coupling called the A -term. However, a simple gauge mediation model predicts $A = 0$ at the messenger scale with which $\mathcal{O}(1\%)$ tuning is necessary (138).

One can try to build a model to generate the A -term. However, such a model requires that the A -term has the same phase as the gaugino masses, since the relative phase is physical and observable. The constraint from item 6 restricts the phase to be smaller than $\mathcal{O}(10^{-(2-3)})$. The same restriction applies for the B parameter in the Higgs potential.

Several ways to address those problems have been considered in the literature. As recent progress in gauge mediation, for example, semi-strongly coupled models are considered in (39) to avoid large hierarchies among $m_{H_u}^2$, $m_{H_d}^2$, $B\mu$ and μ^2 , where the CP phase is controlled by a symmetry. A perturbative model without the hierarchy is presented in (139) where the Higgs and the messenger sectors are extended. (See also (51, 140, 141, 142) for earlier discussions.) In (143), it has been pointed out that the correct size of m_Z can be obtained even in the presence of the hierarchy. An explicit example to realize the correct hierarchy pattern was presented. A mechanism to suppress $m_{H_u}^2$, $m_{H_d}^2$ and $B\mu$ relative to μ^2 by large renormalization effects has been discussed in (144, 145). In (62), discussions were reformulated in the context of general gauge mediation. For another recent approach using more than one supersymmetry breaking spurion, see (146).

One should not forget that natural electroweak symmetry breaking is one of the main motivations to consider the supersymmetric standard model. Supersymmetry breaking and its connection to the Higgs sector should be arranged so that electroweak symmetry breaking happens naturally. We encountered problems in this connection, but, on the other hand, the need for a special structure of the Higgs sector may be a hint for the actual underlying theory. It is desirable to have more ideas to be confirmed in future or on-going experiments, such as the

LHC experiments, in order to reveal the role of supersymmetry in Nature.

Acknowledgments

We would like to thank M. Graesser, Z. Komargodski, and D. Shih for comments on this manuscript. RK is supported in part by the Grant-in-Aid for Scientific Research 21840006 of JSPS. HO is supported in part by US DOE grant DE-FG03-92-ER40701, the World Premier International Research Center Initiative of MEXT of Japan, and by a Grant-in-Aid for Scientific Research (C) 20540256 of JSPS. YO's research at Perimeter Institute for Theoretical Physics is supported in part by the Government of Canada through NSERC and by the Province of Ontario through MRI.

8 NUMBERED LITERATURE CITED

LITERATURE CITED

1. Montonen C, Olive DI, Phys. Lett. B72:117 (1977).
2. Witten E, Olive DI, Phys. Lett. B78:97 (1978).
3. Osborn H, Phys. Lett. B83:321 (1979).
4. Sen A, Phys. Lett. B329:217 (1994), hep-th/9402032.
5. Vafa C, Witten E, Nucl. Phys. B431:3 (1994), hep-th/9408074.
6. Seiberg N, Witten E, Nucl. Phys. B426:19 (1994), hep-th/9407087.
7. Seiberg N, Nucl. Phys. B435:129 (1995), hep-th/9411149.
8. Seiberg N, Phys. Rev. D49:6857 (1994), hep-th/9402044.
9. Maldacena JM, Adv. Theor. Math. Phys. 2:231 (1998), hep-th/9711200.
10. Dimopoulos S, Georgi H, Nucl. Phys. B193:150 (1981).

11. Sakai N, Zeit. Phys. C11:153 (1981).
12. Dimopoulos S, Raby S, Wilczek F, Phys. Rev. D24:1681 (1981).
13. Jungman G, Kamionkowski M, Griest K, Phys. Rept. 267:195 (1996), hep-ph/9506380.
14. Particle Data Group, Amsler C, et al., Phys. Lett. B667:1 (2008).
15. Dine M, Fischler W, Srednicki M, Nucl. Phys. B189:575 (1981).
16. Intriligator KA, Seiberg N, Shih D, JHEP 04:021 (2006), hep-th/0602239.
17. Moroi T, Murayama H, Yamaguchi M, Phys. Lett. B303:289 (1993).
18. Banks T, Kaplan DB, Nelson AE, Phys. Rev. D49:779 (1994), hep-ph/9308292.
19. Bagger J, Poppitz E, Randall L, Nucl. Phys. B426:3 (1994), hep-ph/9405345.
20. Witten E, Nucl. Phys. B188:513 (1981).
21. Dimopoulos S, Raby S, Nucl. Phys. B192:353 (1981).
22. Witten E, Nucl. Phys. B202:253 (1982).
23. Nelson AE, Seiberg N, Nucl. Phys. B416:46 (1994), hep-ph/9309299.
24. Affleck I, Dine M, Seiberg N, Nucl. Phys. B241:493 (1984).
25. Affleck I, Dine M, Seiberg N, Phys. Lett. B137:187 (1984).
26. Affleck I, Dine M, Seiberg N, Nucl. Phys. B256:557 (1985).
27. Izawa KI, Yanagida T, Prog. Theor. Phys. 95:829 (1996), hep-th/9602180.
28. Intriligator KA, Thomas SD, Nucl. Phys. B473:121 (1996), hep-th/9603158.
29. Izawa KI, Takahashi F, Yanagida TT, Yonekura K, Phys. Lett. B677:195 (2009), 0902.3854.
30. Izawa KI, Takahashi F, Yanagida TT, Yonekura K, Phys. Rev. D80:085017 (2009), 0905.1764.

31. Chacko Z, Luty MA, Ponton E, JHEP 12:016 (1998), hep-th/9810253.
32. Murayama H, Phys. Lett. B355:187 (1995), hep-th/9505082.
33. Intriligator KA, Seiberg N, Class. Quant. Grav. 24:S741 (2007), hep-ph/0702069.
34. Polonyi J, Hungary Central Inst Res - KFKI-77-93 (77,REC.JUL 78) 5p.
35. Ooguri H, Ookouchi Y, Park CS, Adv. Theor. Math. Phys. 12:405 (2008), 0704.3613.
36. Marsano J, Ooguri H, Ookouchi Y, Park CS, Nucl. Phys. B798:17 (2008), 0712.3305.
37. Kitano R, Ookouchi Y, Phys. Lett. B675:80 (2009), 0812.0543.
38. Kitano R, Phys. Lett. B641:203 (2006), hep-ph/0607090.
39. Ibe M, Kitano R, JHEP 08:016 (2007), 0705.3686.
40. Murayama H, Nomura Y, Phys. Rev. Lett. 98:151803 (2007), hep-ph/0612186.
41. Murayama H, Nomura Y, Phys. Rev. D75:095011 (2007), hep-ph/0701231.
42. O’Raifeartaigh L, Nucl. Phys. B96:331 (1975).
43. Wess J, Zumino B, Phys. Lett. B49:52 (1974).
44. Ray S, Phys. Lett. B642:137 (2006), hep-th/0607172.
45. Giudice GF, Rattazzi R, Phys. Rept. 322:419 (1999), hep-ph/9801271.
46. Dine M, Fischler W, Phys. Lett. B110:227 (1982).
47. Alvarez-Gaume L, Claudson M, Wise MB, Nucl. Phys. B207:96 (1982).
48. Nappi CR, Ovrut BA, Phys. Lett. B113:175 (1982).
49. Dine M, Nelson AE, Phys. Rev. D48:1277 (1993), hep-ph/9303230.
50. Dine M, Nelson AE, Shirman Y, Phys. Rev. D51:1362 (1995), hep-ph/9408384.

- 51. Dine M, Nelson AE, Nir Y, Shirman Y, Phys. Rev. D53:2658 (1996), hep-ph/9507378.
- 52. Giudice GF, Rattazzi R, Nucl. Phys. B511:25 (1998), hep-ph/9706540.
- 53. Arkani-Hamed N, Giudice GF, Luty MA, Rattazzi R, Phys. Rev. D58:115005 (1998), hep-ph/9803290.
- 54. Givon A, Katz A, Komargodski Z, Shih D, JHEP 10:092 (2008), 0808.2901.
- 55. Intriligator K, Shih D, Sudano M, JHEP 03:106 (2009), 0809.3981.
- 56. Meade P, Seiberg N, Shih D, Prog. Theor. Phys. Suppl. 177:143 (2009), 0801.3278.
- 57. Carpenter LM, Dine M, Festuccia G, Mason JD, Phys. Rev. D79:035002 (2009), 0805.2944.
- 58. Buican M, Meade P, Seiberg N, Shih D, JHEP 03:016 (2009), 0812.3668.
- 59. Nakayama Y, Taki M, Watari T, Yanagida TT, Phys. Lett. B655:58 (2007), 0705.0865.
- 60. Distler J, Robbins D, (2008), 0807.2006.
- 61. Intriligator KA, Sudano M, JHEP 11:008 (2008), 0807.3942.
- 62. Komargodski Z, Seiberg N, JHEP 03:072 (2009), 0812.3900.
- 63. Ooguri H, Ookouchi Y, Park CS, Song J, Nucl. Phys. B808:121 (2009), 0806.4733.
- 64. Buican M, Komargodski Z, (2009), 0909.4824.
- 65. Abel S, Dolan MJ, Jaeckel J, Khoze VV, JHEP 12:001 (2009), 0910.2674.
- 66. Kobayashi T, Nakai Y, Takahashi R, (2009), 0910.3477.
- 67. Seiberg N, Volansky T, Wecht B, JHEP 11:004 (2008), 0809.4437.
- 68. Elvang H, Wecht B, JHEP 06:026 (2009), 0904.4431.
- 69. Ibe M, Izawa KI, Nakai Y, Phys. Rev. D80:035002 (2009), 0907.2970.

- 70. Argurio R, Bertolini M, Ferretti G, Mariotti A, (2009), 0912.0743.
- 71. Dine M, Mason JD, Phys. Rev. D78:055013 (2008), 0712.1355.
- 72. Poppitz E, Trivedi SP, Phys. Rev. D55:5508 (1997), hep-ph/9609529.
- 73. Arkani-Hamed N, March-Russell J, Murayama H, Nucl. Phys. B509:3 (1998), hep-ph/9701286.
- 74. Murayama H, Phys. Rev. Lett. 79:18 (1997), hep-ph/9705271.
- 75. Dimopoulos S, Dvali GR, Rattazzi R, Phys. Lett. B413:336 (1997), hep-ph/9707537.
- 76. Luty MA, Phys. Lett. B414:71 (1997), hep-ph/9706554.
- 77. Csaki C, Shirman Y, Terning J, JHEP 05:099 (2007), hep-ph/0612241.
- 78. Essig R, Fortin JF, Sinha K, Torroba G, Strassler MJ, JHEP 03:043 (2009), 0812.3213.
- 79. Zur BK, Mazzucato L, Oz Y, JHEP 10:099 (2008), 0807.4543.
- 80. Kitano R, Ooguri H, Ookouchi Y, Phys. Rev. D75:045022 (2007), hep-ph/0612139.
- 81. Sato R, Yanagida TT, Yonekura K, (2009), 0910.3790.
- 82. Hall LJ, Randall L, Nucl. Phys. B352:289 (1991).
- 83. Randall L, Rius N, Phys. Lett. B286:299 (1992).
- 84. Dine M, MacIntire D, Phys. Rev. D46:2594 (1992), hep-ph/9205227.
- 85. Fox PJ, Nelson AE, Weiner N, JHEP 08:035 (2002), hep-ph/0206096.
- 86. Amigo SDL, Blechman AE, Fox PJ, Poppitz E, JHEP 01:018 (2009), 0809.1112.
- 87. Benakli K, Goodsell MD, Nucl. Phys. B816:185 (2009), 0811.4409.
- 88. Shih D, JHEP 02:091 (2008), hep-th/0703196.
- 89. Givon A, Katz A, Komargodski Z, JHEP 06:003 (2008), 0804.1805.

- 90. Cheung C, Fitzpatrick AL, Shih D, JHEP 07:054 (2008), 0710.3585.
- 91. Dimopoulos S, Giudice GF, Phys. Lett. B393:72 (1997), hep-ph/9609344.
- 92. Izawa KI, Nomura Y, Tobe K, Yanagida T, Phys. Rev. D56:2886 (1997), hep-ph/9705228.
- 93. Abel S, Durnford C, Jaeckel J, Khoze VV, Phys. Lett. B661:201 (2008), 0707.2958.
- 94. Haba N, Maru N, Phys. Rev. D76:115019 (2007), 0709.2945.
- 95. Sato R, Yonekura K, (2009), 0912.2802.
- 96. Komargodski Z, Shih D, JHEP 04:093 (2009), 0902.0030.
- 97. Heckman JJ, (2010), 1001.0577.
- 98. Giveon A, Kutasov D, Rev. Mod. Phys. 71:983 (1999), hep-th/9802067.
- 99. Ooguri H, Ookouchi Y, Phys. Lett. B641:323 (2006), hep-th/0607183.
- 100. Franco S, Garcia-Etxebarria I, Uranga AM, JHEP 01:085 (2007), hep-th/0607218.
- 101. Bena I, Gorbatov E, Hellerman S, Seiberg N, Shih D, JHEP 11:088 (2006), hep-th/0608157.
- 102. Giveon A, Kutasov D, Nucl. Phys. B778:129 (2007), hep-th/0703135.
- 103. Klebanov IR, Strassler MJ, JHEP 08:052 (2000), hep-th/0007191.
- 104. Kachru S, Pearson J, Verlinde HL, JHEP 06:021 (2002), hep-th/0112197.
- 105. Argurio R, Bertolini M, Franco S, Kachru S, JHEP 01:083 (2007), hep-th/0610212.
- 106. Argurio R, Bertolini M, Franco S, Kachru S, JHEP 06:017 (2007), hep-th/0703236.
- 107. Gherghetta T, Pomarol A, Nucl. Phys. B586:141 (2000), hep-ph/0003129.
- 108. Gherghetta T, Pomarol A, Nucl. Phys. B602:3 (2001), hep-ph/0012378.

109. Goldberger WD, Nomura Y, Tucker-Smith D, Phys. Rev. D67:075021 (2003), hep-ph/0209158.
110. Nomura Y, Tucker-Smith D, Phys. Rev. D68:075003 (2003), hep-ph/0305214.
111. Nomura Y, (2004), hep-ph/0410348.
112. Benini F, et al., JHEP 12:031 (2009), 0903.0619.
113. DeWolfe O, Kachru S, Mulligan M, Phys. Rev. D77:065011 (2008), 0801.1520.
114. McGuirk P, Shiu G, Sumitomo Y, (2009), 0911.0019.
115. McGuirk P, Shiu G, Sumitomo Y, (2009), 0910.4581.
116. Ooguri H, Vafa C, Nucl. Phys. B500:62 (1997), hep-th/9702180.
117. Vafa C, J. Math. Phys. 42:2798 (2001), hep-th/0008142.
118. Maldacena JM, Nunez C, Phys. Rev. Lett. 86:588 (2001), hep-th/0008001.
119. Cachazo F, Intriligator KA, Vafa C, Nucl. Phys. B603:3 (2001), hep-th/0103067.
120. Aganagic M, Beem C, Seo J, Vafa C, Nucl. Phys. B789:382 (2008), hep-th/0610249.
121. Heckman JJ, Seo J, Vafa C, JHEP 07:073 (2007), hep-th/0702077.
122. Aganagic M, Beem C, Freivogel B, Nucl. Phys. B795:291 (2008), 0708.0596.
123. Buican M, Malyshev D, Verlinde H, JHEP 06:108 (2008), 0710.5519.
124. Marsano J, Papadodimas K, Shigemori M, Nucl. Phys. B789:294 (2008), 0705.0983.
125. Marsano J, Papadodimas K, Shigemori M, Nucl. Phys. B804:19 (2008), 0801.2154.
126. LEP Working Group for Higgs boson searches, Barate R, et al., Phys. Lett.

- B565:61 (2003), hep-ex/0306033.
127. DELPHI, Abdallah J, et al., Eur. Phys. J. C31:421 (2003), hep-ex/0311019.
128. Hooper D, Plehn T, Phys. Lett. B562:18 (2003), hep-ph/0212226.
129. Belanger G, Boudjema F, Pukhov A, Rosier-Lees S, (2002), hep-ph/0212227.
130. Dreiner HK, et al., Eur. Phys. J. C62:547 (2009), 0901.3485.
131. Tevatron Electroweak Working Group, (2009), 0903.2503.
132. Baker CA, et al., Phys. Rev. Lett. 97:131801 (2006), hep-ex/0602020.
133. Okada Y, Yamaguchi M, Yanagida T, Prog. Theor. Phys. 85:1 (1991).
134. Ellis JR, Ridolfi G, Zwirner F, Phys. Lett. B257:83 (1991).
135. Haber HE, Hempfling R, Phys. Rev. Lett. 66:1815 (1991).
136. Barbieri R, Giudice GF, Nucl. Phys. B306:63 (1988).
137. Chacko Z, Nomura Y, Tucker-Smith D, Nucl. Phys. B725:207 (2005), hep-ph/0504095.
138. Kitano R, Nomura Y, Phys. Rev. D73:095004 (2006), hep-ph/0602096.
139. Giudice GF, Kim HD, Rattazzi R, Phys. Lett. B660:545 (2008), 0711.4448.
140. Dvali GR, Giudice GF, Pomarol A, Nucl. Phys. B478:31 (1996), hep-ph/9603238.
141. Agashe K, Graesser M, Nucl. Phys. B507:3 (1997), hep-ph/9704206.
142. de Gouvea A, Friedland A, Murayama H, Phys. Rev. D57:5676 (1998), hep-ph/9711264.
143. Csaki C, Falkowski A, Nomura Y, Volansky T, Phys. Rev. Lett. 102:111801 (2009), 0809.4492.
144. Roy TS, Schmaltz M, Phys. Rev. D77:095008 (2008), 0708.3593.
145. Murayama H, Nomura Y, Poland D, Phys. Rev. D77:015005 (2008), 0709.0775.

146. Liu T, Wagner CEM, JHEP 06:073 (2008), 0803.2895.

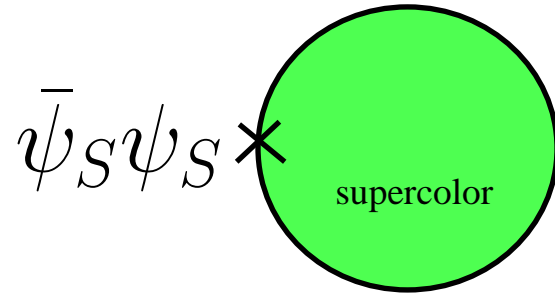


Figure 1: Conceptual picture for dynamical supersymmetry breaking.

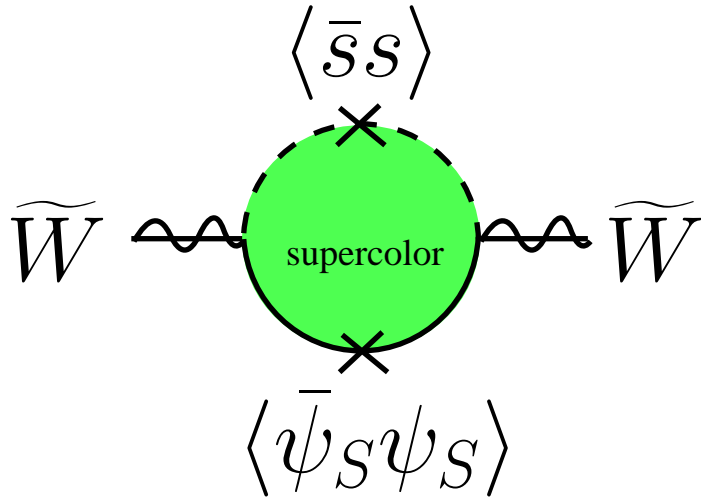


Figure 2: A diagram to contribute gaugino masses.

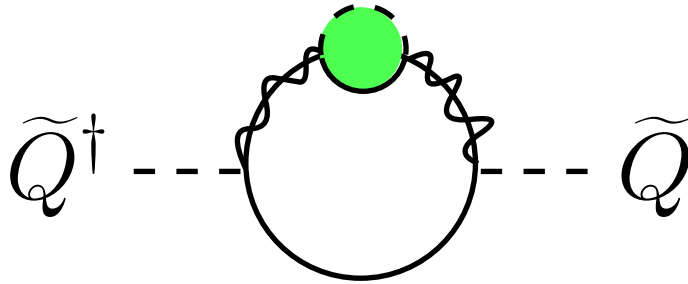


Figure 3: One of the diagrams to contribute sfermion masses.

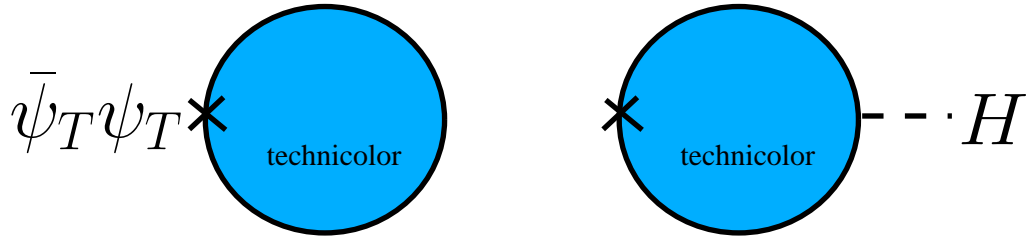


Figure 4: Pictures for electroweak symmetry breaking and the VEV of the Higgs field.

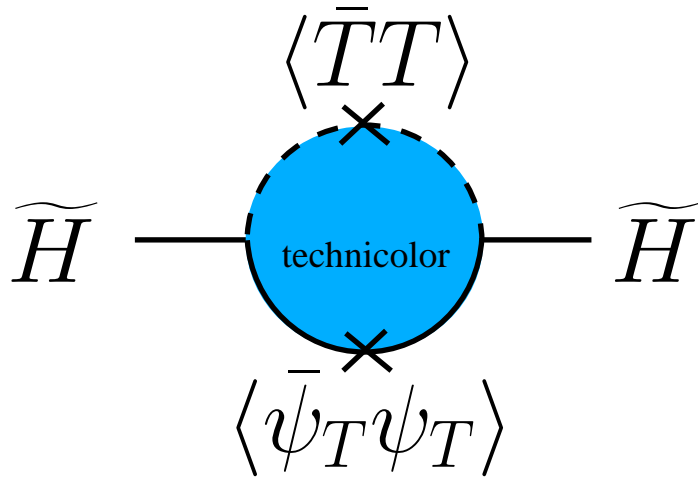


Figure 5: A diagram to contribute the Higgsino mass.

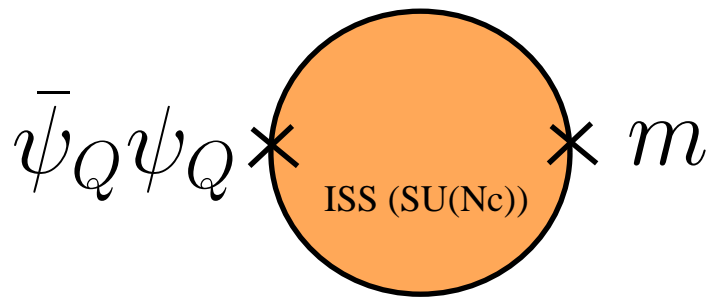


Figure 6: A picture for dynamical supersymmetry breaking in the ISS model.

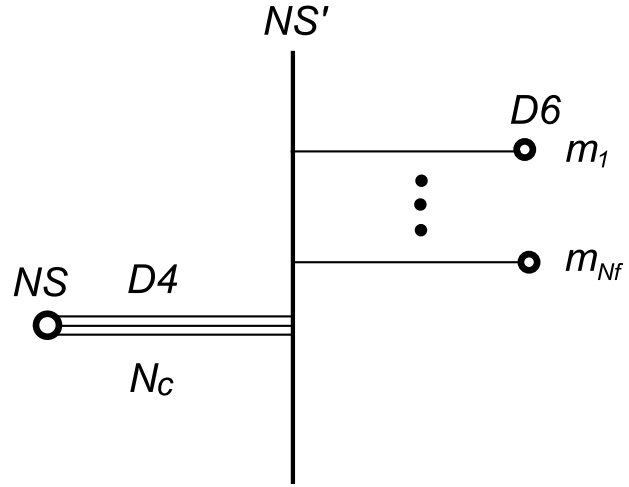


Figure 7: The electric brane configuration for $\mathcal{N} = 1$ supersymmetric gauge theory with massive flavors.

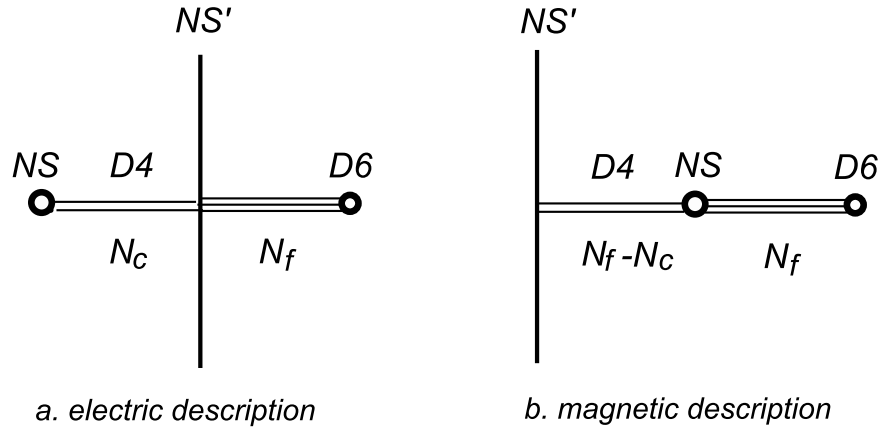


Figure 8: The supersymmetric brane configurations for $\mathcal{N} = 1$ supersymmetric gauge theory with massless flavors.

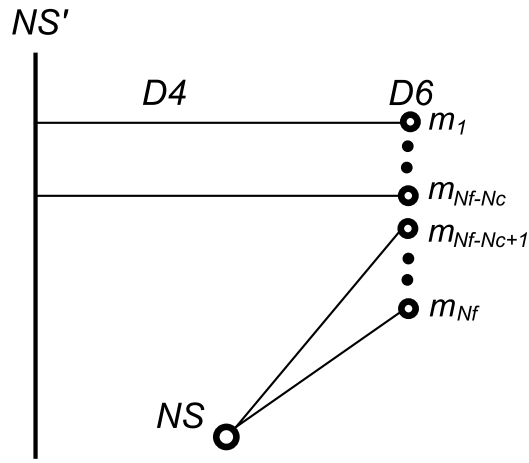


Figure 9: The brane configuration for the meta-stable supersymmetry breaking vacuum.